Quantum-Limited Sensitivity of Single-Electron-Transistor-Based Displacement Detectors

D. Mozyrsky, I. Martin, and M. B. Hastings
Theoretical Division, Los Alamos National Laboratory, Los Alamos, New Mexico 87545, USA
(Received 18 June 2003; published 8 January 2004)

We consider a model of a quantum-mechanical resonator capacitively coupled to a single electron transistor (SET). The tunnel current in the SET is modulated by the vibrations of the resonator, and thus the system operates as a displacement detector. We analyze the effect of the backaction noise of charge fluctuations in the SET onto the dynamics of the resonator and evaluate the displacement sensitivity of the system. The relation between the "classical" and "quantum" parts of the SET charge noise and their effect on the measured system are also discussed.

DOI: 10.1103/PhysRevLett.92.018303 PACS numbers: 85.85.+j, 05.30.–d, 05.60.Gg

Micromechanical resonators have been used as ultrasensitive force detectors in a number of experimental applications, ranging from atomic force microscopy to magnetic resonance force microscopy [1] to experiments on Casimir force detection [2]. Recently, mechanical resonators with vibrational eigenfrequencies (ν) of the order of 1 GHz have been fabricated [3]. At low temperatures ($h\nu > k_B T \sim 50$ mK), these resonators provide an example of a man-made system that can be used to test the basic principles of quantum mechanics at the macroscopic level [4].

The standard cantilever displacement measurement schemes are based on laser interferometry, and can reach the levels of sensitivity of the order 10^{-4} Å$/\sqrt{\text{Hz}}$. This level of sensitivity requires, however, high laser power that may not be compatible with the ultralow temperature operation. This limitation provided the motivation to explore alternative *electrical* measurement schemes [5–7]. In particular, Blencowe and Wybourne [5] have suggested based on a semiclassical analysis that by capacitively coupling the cantilever to a single electron transistor (SET), it is possible to achieve the sensitivity better than the zero-point-motion uncertainty. More recently, two of us [7] have found, based on a fully quantum-mechanical description of the quantum measurement of a cantilever using a quantum point contact (QPC), that the apparatus backaction (current shot noise that induces force noise on the cantilever) fundamentally limits the displacement sensitivity and leads to a quantum-to-classical transition in the oscillator dynamics. Because of the resonant nature of transport through SET, it is expected to have significantly higher displacement sensitivity than the QPC, and hence is more attractive from the experimental standpoint. Here, we analyze the fundamental sensitivity limits of an SET-based detection scheme. We find that the higher classical sensitivity comes at the expense of drastically increased back action (island charge noise), which leads to large rms fluctuations of the cantilever and to deterioration of the oscillator quality factor in the experimentally most attractive *threshold* regime. The optimal sensitivity is achieved in the *cotunneling* regime in which the SET scheme becomes equivalent to the QPC-based detection. Similar conclusions have also been reached in the study of SET charge sensitivity [8]. We find that under no circumstances is it possible to exceed the standard quantum limit using continuous SET-based detection.

The model is schematically presented in Fig. 1. The measuring apparatus is a single electron transistor: a quantum dot coupled to the leads via tunnel junctions. For simplicity we assume that the dot contains a single resonant level. The mechanical system, which we refer to as an oscillator and which can be either a micromechanical resonator or a localized phonon mode, is capacitively coupled to the resonant level. Then, the oscillator displacement can alter the position of the resonant level with respect to the chemical potentials in the leads and thus affect the tunnel current through the device.

The Hamiltonian of the model can be written as

$$H = H_{\text{leads}} + H_{\text{osc}} + H', \tag{1}$$

where the first two terms are the Hamiltonians of the electrons in the leads and the oscillator, respectively, $H_{\text{leads}} = \sum_{q,n=L,R} \epsilon_{qn} c^{\dagger}_{qn} c_{qn}$ and $H_{\text{osc}} = (-1/2m)\partial_x^2 + V(x)$. Here $c^{\dagger}_{q,L(R)}$ ($c_{q,L(R)}$) creates (annihilates) an electron with a quantum number q in the left (right) lead, m and ω_0 are the oscillator effective mass and frequency,

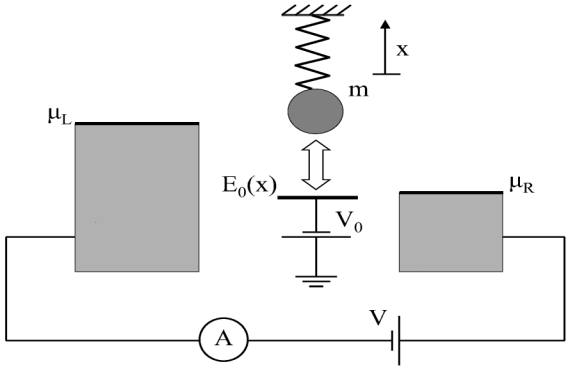

FIG. 1. Schematics for the model setup.

and x is the oscillator coordinate. Here and in the following we set both Planck's constant $\hbar$ and the electron charge e equal to unity, unless stated otherwise. The single particle states ϵ_{qn} are filled up to chemical potentials in the leads, μ_L and μ_R, which are biased by external voltage, $\mu_L - \mu_R = V$; see Fig. 1. For simplicity we assume zero external temperature. The Hamiltonian H' includes both electron tunneling and modulation of the position of the resonant level by the oscillator:

$$H' = \sum_{q,n=L,R} T_n (d^\dagger c_{qn} + c_{qn}^\dagger d) + E_0(x)\hat{n}_0. \tag{2}$$

In the first, tunneling, term of the Hamiltonian (2), the operator $d_0^\dagger$ (d) creates (annihilates) an electron in the resonant level, $\hat{n}_0 = d^\dagger d$, and the tunneling amplitude T_n is assumed to be independent of the single particle states in the leads. We assume that the energy of the resonant level $E_0(x)$ depends linearly on the oscillator's coordinate x, i.e., $E_0(x) = \epsilon_0 + \lambda x$. The unperturbed position of the resonant level ϵ_0 can be varied by appropriately adjusting the gate voltage V_0; see Fig. 1. The parameter λ physically represents an effective electric field in the capacitor formed by the oscillator and the quantum dot.

We use the Keldysh-Feynman-Vernon formalism [9,10] to determine the evolution of the oscillator under the influence of the tunneling electrons. We define a scattering operator for the oscillator alone, i.e., with electronic degrees of freedom traced out:

$$S_{\rm osc} = {\rm Tr}_{\rm el}[\rho_{\rm el} \mathcal{T}_{\rm c}\, S(-\infty, \infty) S(\infty, -\infty)]/{\rm Tr}[\rho_{\rm el}]. \tag{3}$$

In Eq. (3) $S(\infty, -\infty)$ and $S(-\infty, \infty)$ are scattering operators for the full system, $S(\infty, -\infty) = \exp[-i \int_{-\infty}^{\infty} H dt]$, where H is defined in Eqs. (1) and (2), and the operator $\mathcal{T}_{\rm c}$ denotes time ordering along the Keldysh contour. The density matrix of the unperturbed electrons is the direct product of the uncoupled density matrices of electron reservoirs in the leads (ρ_L and ρ_R) with an empty electron state in the resonant level ($\rho_D = dd^\dagger$), $\rho_{\rm el} = \rho_L \otimes \rho_R \otimes \rho_D$. Equation (3) implies that at $t = -\infty$ the leads, the resonant level, and the oscillator are uncoupled and that the interaction, H', is switched on adiabatically at $t > -\infty$.

In what follows we assume that the coupling constant λ is small, while the tunneling amplitudes T_L and T_R need not be small. Then, $S_{\rm osc}$ can be written explicitly as a functional integral over the oscillator coordinate as

$$S_{\rm osc} = \int \mathcal{D}x \exp\left[i \int_{\rm c} dt \mathcal{L}'_{\rm osc} \right] \times \exp\left[-\frac{\lambda^2}{2} \int_{\rm c} dt_1 dt_2 x(t_1) x(t_2) K(t_1 - t_2) + \ldots \right], \tag{4}$$

where the contour is closed at $t = +\infty$ but not at $t = -\infty$, and where we consider only $O(\lambda^2)$ contributions to $S_{\rm osc}$. The higher orders in the expansion [denoted by ... in Eq. (4)] are unimportant in the limit of strong tunneling as will be seen below.

The first order contribution in λ, i.e., interaction of the oscillator with average charge $\langle \hat{n}_0 \rangle$ in the dot, is included in the Lagrangian of the oscillator in Eq. (4), $\mathcal{L}'_{\rm osc} = \mathcal{L}^{\rm (bare)}_{\rm osc} - \lambda \langle \hat{n}_0 \rangle_{\rm el} x$. The charge $\langle \hat{n}_0 \rangle$ is related to the Fourier transform of the *renormalized* single particle Green's function $G_D(t' - t) = -i \langle \mathcal{T}_{\rm c} d(t) d^\dagger(t') \rangle_{\rm el}$ as $\langle n_0 \rangle_{\rm el} = (1/2\pi i) \int d\omega G_D^{-+}(\omega)$. The averaging denoted by $\langle \rangle_{\rm el}$ is taken with respect to the exact stationary state of the electronic subsystem alone, i.e., decoupled from the oscillator. The renormalization of $G_D(t)$ by the tunneling transitions can be obtained by a standard calculation [11]. Following the notation of Refs. [9,12] we define a matrix Green's function $G_D^{ij}(t_2 - t_1) = -i \langle d(t_1^j) d^\dagger(t_2^i) \rangle_{\rm el}$, where t_1^i and t_2^j can be on either the same or different Keldysh contours, i.e., $i, j = \pm$. We also introduce the unperturbed Green's functions of the electrons in the left and the right leads $G_{qn,0}^{ij}(t_1 - t_2) = -i \langle \mathcal{T}_{\rm c}\, c(t_1^j)_{qn} c^\dagger(t_2^i)_{qn} \rangle_{{\rm el},0}$, $n = L, R$, and the unperturbed Green's function of the dot electron $G_{D,0}^{ij}(t_1 - t_2) = -i \langle \mathcal{T}_{\rm c}\, d(t_1^j) d^\dagger(t_2^i) \rangle_{{\rm el},0}$. The time ordered and antitime ordered Green's functions, i.e., with time arguments on forward and return branches, respectively, can be expressed in terms of the Green's functions with time arguments on different branches as $G^{++}(t) = \Theta(t) G^{-+}(t) + \Theta(-t) G^{+-}(t)$ and $G^{--}(t) = \Theta(t) G^{-+}(t) + \Theta(-t) \times G^{+-}(t)$, where $\Theta(t)$ is a unit step function [9,12]. Then, by solving the Dyson equation $G_D^{ij} = G_{D,0}^{ij} + G_{D,0}^{ik} \Sigma^{kl} G_D^{lj}$, where self-energy $\Sigma^{ij} = \sum_{q,n=L,R} T_n^2 G_{qn,0}^{ij}$, after straightforward calculation we obtain

$$G_D^{-+}(\omega) = 2i \frac{\Gamma_L \Theta(\mu_L - \omega) + \Gamma_R \Theta(\mu_R - \omega)}{(\omega - \epsilon_0)^2 + (\Gamma_L + \Gamma_R)^2}, \tag{5a}$$

$$G_D^{+-}(\omega) = -2i \frac{\Gamma_L \Theta(\omega - \mu_L) + \Gamma_R \Theta(\omega - \mu_R)}{(\omega - \epsilon_0)^2 + (\Gamma_L + \Gamma_R)^2}, \tag{5b}$$

where $G_D^{ij}(\omega) = \int G_D^{ij}(t) \exp(i\omega t) dt$. In Eq. (5) we introduced tunneling rates $\Gamma_{L(R)} = \pi T_{L(R)}^2 \rho_{L(R)}$, where the density of states in the leads $\rho_{L(R)}$ are assumed constant for simplicity.

The $O(\lambda^2)$ contribution to the effective action in Eq. (4) is generated by the integral kernel $K(t_1 - t_2)$, which is just a two-point correlation function of charge fluctuations in the dot, and can be expressed as a product of two single particle Green's functions G_D. The double integral in Eq. (4) can be rewritten as

$$\int_{-\infty}^{\infty} dt_1 dt_2 \{ 2i x^c(t_1) x^q(t_2) \Theta(t_1 - t_2) A(t_1 - t_2) + x^q(t_1) x^q(t_2) S(t_1 - t_2) \}, \tag{6}$$

where we have introduced the "rotated" Keldysh variables $x^q(t) = x(t^+) - x(t^-)$, $x^c(t) = x(t^+) + x(t^-)$. The kernels $A(t_1 - t_2) = {\rm Im}[G_D^{+-}(t_2 - t_1) G_D^{-+}(t_1 - t_2)]$ and $S(t_1 - t_2) = {\rm Re}[G_D^{+-}(t_2 - t_1) G_D^{-+}(t_1 - t_2)]$, with G_D^{+-} and G_D^{-+} given by Eqs. (4), are related to antisymmetric

(quantum) and symmetric (classical) parts of the charge correlation function $K(t_2 - t_1)$.

Suppose that the kernel $K(t_2 - t_1)$ is nonzero on a time scale which is much shorter than the oscillator period. Then, we can approximate $x^{c(q)}(t_2) \simeq x^{c(q)}(t_1) + \dot{x}^{c(q)}(t_1)(t_1 - t_2)$ in Eq. (6), which yields an action local in time:

$$\mathcal{F} = -\frac{\lambda^2}{2}\int dt(iRx^c x^q + iA\dot{x}^c x^q + Sx^q x^q) + \ldots, \quad (7)$$

where the coefficients R, A, and S can be expressed in terms of single particle Green's functions as follows:

$$S = \frac{1}{2\pi}\int d\omega G_D^{-+}(\omega)G_D^{+-}(\omega), \quad (8a)$$

$$A = \frac{1}{2\pi}\int d\omega G_D^{-+}(\omega)\frac{\partial}{\partial\omega}G_D^{+-}(\omega), \quad (8b)$$

$$R = \frac{1}{\pi}\mathcal{P}\int d\omega_1 d\omega_2 \frac{G_D^{-+}(\omega_1)G_D^{+-}(\omega_2)}{\omega_1 - \omega_2}. \quad (8c)$$

The effective action in Eq. (7) is exactly of the form of the Caldeira and Leggett action [13], derived for a bosonic heat bath at high temperature. The first term in the effective action $\mathcal{F}$ is a renormalization of the oscillator potential. In contrast to the infinite renormalization in the Caldeira-Leggett model, here R is finite. Evaluating the integral in Eq. (8c) using Eqs. (5) yields $R = 2\partial\langle n_0\rangle_{\rm el}/\partial\epsilon_0$. For a linear oscillator, $V(x) = m\omega_0^2 x^2/2$, this term provides a frequency shift $\Delta\omega_0 = -\lambda^2 R/(2m\omega_0)$. More generally, the R term in Eq. (7) can be combined with the first order renormalization in Eq. (4) to yield the effective potential of the oscillator as $V(x) \simeq V^{\rm bare}(x) + \lambda x\langle n_0(x)\rangle_{\rm el}$, where $\langle n_0(x)\rangle_{\rm el}$ is the occupation number of the resonant level for a fixed position of the oscillator. The last term in Eq. (7) provides dephasing of the system. Its effect at the classical level corresponds to a white noise force $f(t)$ exerted by the tunnel current on the oscillator. The second term causes energy damping. The classical equation of motion for the oscillator [13] can be written as $m\ddot{x} + m\gamma\dot{x} + \partial_x V = f(t)$, where, in our case, $m\gamma = \lambda^2 A$ and $\langle f(t)f(t')\rangle = \lambda^2 S\delta(t - t')$. Thus the classical and quantum parts of the resonant level charge correlation function determine fluctuations and dissipation for the oscillator, respectively. One can therefore define an effective temperature, $T_{\rm eff}$, using a fluctuation-dissipation relation, giving $S/A = 2T_{\rm eff}$. The effective temperature, $T_{\rm eff}$, is not determined by the reservoir's actual temperature, as in the Caldeira-Leggett model [13], but rather by the coupling to the tunnel current.

$T_{\rm eff}$ determines the fluctuations of the oscillator coordinate due to the tunnel current induced noise. In the case of a linear oscillator the dispersion of the oscillator coordinate is $\langle x^2\rangle = T_{\rm eff}/(m\omega_0^2)$. We can now check the validity of our expansion in λ. From the structure of Eqs. (7) and (8), we see that higher order terms in Eq. (4) will be smaller by powers of the dimensionless parameter $\lambda\sqrt{\langle x^2\rangle}/(\Gamma_L + \Gamma_R)$. Physically, if the oscillator induced shift of the resonant level is small compared to the width of the level, the backaction on the oscillator is weakly dependent on the position of the oscillator and the higher order nonlinearities are unimportant. Thus, for sufficiently large Γ, we need to consider only the leading, quadratic, terms in the effective action (7).

It is instructive to evaluate $T_{\rm eff}$ explicitly, using Eqs. (5) and (8) in two limiting cases—the threshold and cotunneling regimes. Suppose first that the resonant level is in the vicinity of one of the chemical potentials in the leads, say, $\epsilon_0 = \mu_R = 0$, and the bias between the chemical potentials is large, $\mu_L = \infty$. This threshold regime corresponds to the situation when the gradient of the current through the SET is steep with respect to the gate voltage V_0, and therefore one can expect that the current is quite sensitive to the displacement of the cantilever [5]. In this case we obtain $\gamma^{\rm thr} = \hbar\lambda^2\Gamma_R/(\pi m\Gamma^3)$, $T_{\rm eff}^{\rm thr} = \pi\Gamma_L/4$. Therefore, the effective temperature of the oscillator in the threshold regime is essentially defined by the tunneling induced width Γ of the resonant level. For a practical Si nanomechanical resonator with dimensions 3 μm $\times$ 0.1 μm $\times$ 0.1 μm ($f_0 \simeq 100$ MHz), coupled to SET with $\Gamma \sim 10^9$ s^{-1} by an effective electric field in the cantilever-dot capacitor $\lambda \sim 10^7$ V/m, this corresponds to an effective temperature $T_{\rm eff} \sim 0.1$ K, and the damping coefficient $\gamma \sim 10^7$ s^{-1}. The effect of backaction in this case limits the lowest achievable oscillator temperature to 0.1 K, and the maximum quality factor to about 100.

Another important limiting regime corresponds to the situation when the resonant level is far above or below the chemical potentials in the leads. This is the so-called cotunneling regime. The tunneling electrons can now occupy the level only virtually and the effective coupling constant between the leads is small as it is suppressed by the large energy separation between the chemical potentials in the leads and the resonant level. Assuming $\mu_L - \mu_R \equiv V \ll (\mu_L + \mu_R)/2 \equiv \mu_F$ and $\mu_{L(R)} \gg \Gamma, \epsilon_0$, we obtain $\gamma^{\rm cot} = \hbar\lambda^2\Gamma^2/(\pi m\mu_F^4)$, $T_{\rm eff}^{\rm cot} = \Gamma_L\Gamma_R V/\Gamma^2$. Compared to the threshold regime, in the cotunneling both the oscillator damping and the backaction noise are significantly reduced, with the bias voltage across the leads determining the effective temperature of the oscillator. This result is consistent with Refs. [7,14].

We are now in a position to analyze the sensitivity of the system. Suppose that the oscillator (which will be assumed linear from now on) is perturbed by an external force $F(t)$, say, a short kick of duration $\tau_F \ll \omega_0^{-1}$, so that $F(t) \simeq F\tau_F\delta(t)$. This kick results in the variation of the oscillator's amplitude by the amount $\delta x = F\tau_F/(m\omega_0)$. What minimum δx can be detected by observing the tunnel current, given the noise $\langle|I_\omega|^2\rangle$ in the current?

The ability to measure a signal can be represented by the integrated signal-to-noise ratio [15]

$$s/n = (1/2\pi)\int d\omega|\langle S_F(\omega)\rangle|^2/\langle|I_\omega|^2\rangle, \quad (9)$$

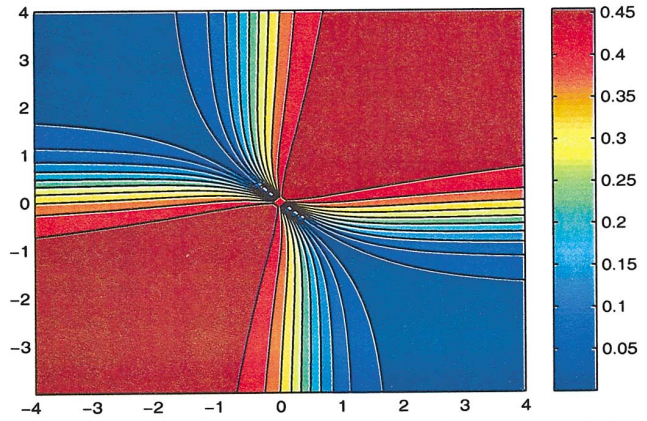

FIG. 2 (color). Sensitivity in dimensionless units ($x_0^2/\delta x_{\min}^2$) as a function of chemical potentials for a symmetric structure ($\Gamma_L = \Gamma_R = \Gamma/2$). The axes are μ_L/Γ and μ_R/Γ.

where $S_F(\omega)$ is the Fourier transform of the detector's response to an external perturbation, i.e., the force F. The variation of the current through the structure due to the variation of the oscillator coordinate is $\delta I = (\partial I/\partial x)\delta x$, and therefore the response can be written as $S_F(\omega) = (\partial I/\partial x)F(\omega)/[m(\omega^2 + i\gamma\omega - \omega_0^2)]$.

The current and the noise can easily be evaluated if we recall that the dynamics of the oscillator is slow compared to the dynamics of the tunneling electrons, i.e., $\omega_0 \ll \Gamma$. The current in the adiabatic approximation, i.e., for a fixed position of the oscillator is given by $I(x) = (1/2\pi)\int_{\mu_R}^{\mu_L} d\omega T(\omega, x)$, where the transmission coefficient $T(\omega, x) = 4\Gamma_L\Gamma_R/\{[\omega - E_0(x)]^2 + \Gamma^2\}$ [11]. For small coupling between the SET and the oscillator, i.e., for $\lambda^2\langle x^2\rangle \ll \Gamma^2$, the noise at low frequencies (of order ω_0) is given by $\langle|I_\omega|^2\rangle \simeq \langle|I_0|^2\rangle + (\partial I/\partial x)^2\langle|\Delta x_\omega|^2\rangle$, where the shot noise is related to the transmission coefficient as $\langle|I_0|^2\rangle = (1/2\pi)\int_{\mu_R}^{\mu_L} d\omega T(\omega, 0)[1 - T(\omega, 0)]$ [16], and $\langle|\Delta x_\omega|^2\rangle = 2\gamma T_{\rm eff}/m[(\omega^2 - \omega_0^2)^2 + \gamma^2\omega^2]$ is the fluctuation spectrum for the oscillator. To find the minimum force detectable by our apparatus, we substitute these expressions into Eq. (9) and set $s/n = 1$; expressing this force in terms of the resulting displacement $\delta x_{\min}$ gives

$$x_0^2/\delta x_{\min}^2 \simeq |\partial I/\partial \epsilon_0|/(4\sqrt{S\langle|I_0|^2\rangle}), \qquad (10)$$

where $x_0^2 = \hbar/(2m\omega_0)$ is amplitude of the zero point motion for the oscillator, and S is given by Eq. (8).

In Fig. 2 we present the sensitivity (10) as a function of chemical potentials in the leads relative to the position of the resonant level. The sensitivity is maximal in the cotunneling regime, where it reaches 1/2. In the threshold regime the sensitivity is somewhat smaller ($\simeq$ by a factor of 2). The sensitivity is worse when the resonant level is positioned symmetrically with respect to the chemical potentials (the blue regions in Fig. 2), as the current sensitivity $\partial I/\partial x$ vanishes in this regime. These results can be compared to conclusions reached in studies of the "nonideality" of SET detectors [8], with the role of quantum noise being an essential difference in our work. Qualitatively, the sensitivity reduction in the threshold tunneling regime is due to the detector latency during the electron dwell time on the island, which contributes to the backaction noise, but not to the measurement.

In summary, we have analyzed the quantum measurement of a mechanical oscillator coupled to an electronic resonant level that models a single electron transistor. We determined the backaction effects of the detector on the quantum system, and found the fundamental sensitivity limits of the scheme in all operation regimes.

We acknowledge useful discussions with K. Schwab and D. Pelekhov. This work was supported by DARPA MOSAIC program and U.S. DOE (D. M. and I. M.), NSF ECS-0102500 (D. M.), DMR-0121146 (D. M.), and U.S. DOE W-7405-ENG-36 (M. B. H.).

[1] D. Rugar, C. S. Yannoni, and J. A. Sidles, Nature (London) **360**, 563 (1992).
[2] U. Mohideen and A. Roy, Phys. Rev. Lett. **81**, 4549 (1998); E. Buks and M. L. Roukes, Phys. Rev. B **63**, 033402 (2001).
[3] X. M. Henry Huang, C. A. Zorman, M. Mehregany, and M. L. Roukes, Nature (London) **421**, 496 (2003).
[4] A. Cho, Science **299**, 36 (2003); A. D. Armour, M. P. Blencowe, and K. C. Schwab, Phys. Rev. Lett. **88**, 148301 (2002).
[5] M. P. Blencowe and M. N. Wybourne, Appl. Phys. Lett. **77**, 3845 (2000); Y. Zhang and M. P. Blencowe, J. Appl. Phys. **91**, 4249 (2002).
[6] R. Knobel and A. N. Cleland, Appl. Phys. Lett. **81**, 2258 (2002); A. N. Cleland *et al.*, *ibid.* **81**, 1699 (2002).
[7] D. Mozyrsky and I. Martin, Phys. Rev. Lett. **89**, 018301 (2002).
[8] D. V. Averin, in *Scalable Quantum Computers*, edited by S. L. Braunstein and H.-K. Lo (Wiley-VCH, Berlin, 2001); A. Maassen van den Brink, Europhys. Lett. **58**, 562 (2002).
[9] L. V. Keldysh, Sov. Phys. JETP **20**, 1018 (1965).
[10] R. P. Feynman and F. L. Vernon, Ann. Phys. (N.Y.) **24**, 118 (1963).
[11] Y. Meir and N. S. Wingreen, Phys. Rev. Lett. **68**, 2512 (1992); A.-P. Jauho, N. S. Wingreen, and Y. Meir, Phys. Rev. B **50**, 5528 (1994).
[12] J. Rammer and H. Smith, Rev. Mod. Phys. **58**, 323 (1986).
[13] A. O. Caldeira and A. J. Leggett, Physica (Amsterdam) **121A**, 587 (1983).
[14] M. B. Hastings, I. Martin, and D. Mozyrsky, Phys. Rev. B **68**, 035101 (2003).
[15] V. B. Braginsky and F. Ya. Khalili, *Quantum Measurement* (Cambridge University Press, Cambridge, U.K., 1992).
[16] V. A. Khlus, Sov. Phys. JETP **66**, 1243 (1987); G. B. Lesovik, JETP Lett. **49**, 592 (1989).